

Study of the anomalous cross section lineshape of $e^+e^- \rightarrow D\bar{D}$ at $\psi(3770)$ with an effective field theory

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We study the anomalous cross section lineshape of $e^+e^- \rightarrow D\bar{D}$ with an effective field theory. Near threshold, most of the $D\bar{D}$ pairs are from the decay of $\psi(3770)$. Taking into account that the non-resonance background is dominated by the $\psi(2S)$ transition, the produced $D\bar{D}$ pair can undergo final state interactions before it is detected. We propose an effective field theory for the low energy $D\bar{D}$ interactions to describe this final state interactions, and find that the anomalous lineshape of the $D\bar{D}$ cross section observed by the BESII collaboration can be well described.

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As the first charmonium state above the $D\bar{D}$ threshold, the resonance $\psi(3770)$ is different from other charmonia with lower masses. Since the $\psi(3770)$ decay into the open charm $D\bar{D}$ is allowed by the OZI rule, this dominant decay mode leads to a broad width up to 27.2 ± 1.0 MeV [1]. Obviously, the direct production process of $e^+e^- \rightarrow \psi(3770) \rightarrow D\bar{D}$ is useful for the study of the property of $\psi(3770)$. In Ref. [2], BESII collaboration reported an anomalous behaviour of the cross section lineshape at the $\psi(3770)$ mass region in $e^+e^- \rightarrow D\bar{D}$ that cannot be described by a simple Breit-Wigner of $\psi(3770)$. Such an observation has inspired interesting theoretical discussions [3–6]. In particular, it was found that the interfering effect between $\psi(3770)$ and $\psi(2S)$ plays a very important role in understanding the anomalous lineshape of $D\bar{D}$ at the $\psi(3770)$ resonance [3, 4]. Such an interference can be recognized by a relative phase factor $e^{i\phi}$ which is introduced between these two resonances, and the phase angle ϕ turns out to be large in order to describe the anomalous $D\bar{D}$ lineshape.

In principle, the phase factor $e^{i\phi}$ can come from the final state interactions of $D\bar{D}$. Thus, it should be interesting to study the $D\bar{D}$ anomalous lineshape using the effective field theory to describe the $D\bar{D}$ final state interactions. This forms our motivation for this work. Near threshold, the $D\bar{D}$ pair produced in $e^+e^- \rightarrow D\bar{D}$ comes from the decay of the $\psi(3770)$ and other non-resonance background processes. Once the $D\bar{D}$ pair was produced, it could undergo final state interactions before it converts into the final observed $D\bar{D}$ state. This could explain the relative phase between the $\psi(3770)$ and other non- $\psi(3770)$ amplitude, and provide a description of the $D\bar{D}$ lineshape. We note that there are several cases in which the final state interactions play important roles in the understanding of the cross section lineshapes [7–10].

It is well-known that effective field theory is a useful tool to study the low energy hadron interactions. An effective field theory makes use of the Taylor expansion on the small ratio between the typical small scale p and the cutoff scale Λ . For example, in Chiral Perturbation Theory (ChPT), p is the momentum of the low energy pion or pion mass while $\Lambda = M_{\rho(770)}$ sets the cutoff scale of this effective theory. An effective field theory for the low energy $D\bar{D}$ is different from that for the low energy $\pi\pi$ interaction because the $\psi(3770)$ should be included explicitly into the effective Lagrangian. Besides the three-vector momentum of the $D(\bar{D})$ meson, another small scale $\delta = M_{\psi(3770)} - 2M_D \approx 40$ MeV also appears in the effective theory. This additional small scale will make the power counting to be different from that in ChPT. Systematic development of the effective field theory with resonances as intermediate states is still under exploration, and interesting discussions on this subject can be found in Refs. [11, 12].

In this work, we use the effective field theory to study the $D\bar{D}$ interaction in order to understand the dynamic details of the anomalous cross section lineshape observed by the BESII Collaboration [2].

At the beginning, we assume that the production of $D\bar{D}$ in e^+e^- annihilation can be approximated by the vector meson dominance (VMD). It means that the cross section for $e^+e^- \rightarrow D\bar{D}$ is dominated by intermediate vector meson productions via $e^+e^- \rightarrow \gamma^* \rightarrow \mathcal{R}_i \rightarrow D\bar{D}$, where \mathcal{R}_i denotes any vector meson with isospin $I = 0$ or $I = 1$. However, it is impossible to sum the contributions from all of \mathcal{R}_i in reality. As a reasonable approximation, one can include the contributions from the vector mesons in the vicinity of the considered energy region, but neglect those far off-shell vector mesons. In the energy region of

the BES data from 3.74 GeV to 3.8 GeV, one can expect that $\psi(3770)$ plays the most important role among all of R_i , while the contributions from all the other R_i can be treated as background. As shown in Ref. [3], the contribution from $\psi(2S)$ dominates the background, while contributions from other states are negligible. Therefore, we only include the contributions from the resonances $\psi(3770)$ and $\psi(2S)$, and neglect those from the other resonances. Namely, $\psi(2S)$ would be the main background near the $D\bar{D}$ threshold.

In VMD [13, 14], the coupling between the vector meson and virtual photon can be described as

$$\mathcal{L}_{V\gamma} = \frac{eM_V^2}{f_V} V_\mu A^\mu, \quad (1)$$

where V_μ is the vector meson field, A_μ is the photon field, and M_V is the mass of the vector meson. Setting the electron mass $m_e \approx 0$, the coupling can be obtained as

$$\frac{e}{f_V} = \left[\frac{3\Gamma_{ee}}{\alpha M_V} \right]^{1/2}, \quad (2)$$

where Γ_{ee} is the electron-position decay width of V_μ , and $\alpha = 1/137$ is the fine-structure constant.

Once the $D\bar{D}$ pair is produced from the decay of the vector meson $\psi(3770)$ or $\psi(2S)$, it can undergo final state interactions through the rescattering processes $D\bar{D} \rightarrow D\bar{D} \rightarrow \cdots \rightarrow D\bar{D}$ which can be described by the effective field theory. In the energy region which is concerned, the three-vector momentum p of the $D(\bar{D})$ meson is small. Thus, it is possible to construct an effective field theory for the low energy $D\bar{D}$ interactions making use of the expansion on the small momentum p . Since the mass of $\psi(3770)$ is just above the threshold of $D\bar{D}$, we need to include $\psi(3770)$ explicitly in the formulation. Near threshold, the $D(\bar{D})$ meson can be treated as nonrelativistic. Thus, the interaction Lagrangian for the $D\bar{D}$ system with the quantum number $J^{PC} = 1^{--}$ can be constructed as

$$\begin{aligned} \delta\mathcal{L} &= \mathcal{L}_{\psi D\bar{D}} + \mathcal{L}_{(D\bar{D})^2} \\ \mathcal{L}_{\psi D\bar{D}} &= ig_{\psi D\bar{D}} \{D^\dagger \nabla \bar{D} - \nabla D^\dagger \bar{D}\} \cdot \psi + ig_{\psi D\bar{D}} \{\bar{D}^\dagger \nabla D - \nabla \bar{D}^\dagger D\} \cdot \psi, \\ \mathcal{L}_{(D\bar{D})^2} &= f_1 \{D^\dagger \nabla \bar{D} - \nabla D^\dagger \bar{D}\} \cdot \{\nabla \bar{D}^\dagger D - \bar{D}^\dagger \nabla D\} \\ &\quad + f_3 \{D^\dagger \tau^i \nabla \bar{D} - \nabla D^\dagger \tau^i \bar{D}\} \cdot \{\nabla \bar{D}^\dagger \tau^i D - \bar{D}^\dagger \tau^i \nabla D\} + \cdots, \\ \text{with } D &= \begin{pmatrix} D^0 \\ D^+ \end{pmatrix}, \quad \bar{D} = \begin{pmatrix} \bar{D}^0 \\ D^- \end{pmatrix}. \end{aligned} \quad (3)$$

Here ψ is the field operator of $\psi(3770)$, $D(\bar{D}^\dagger)$ annihilates a $D(\bar{D})$ meson while $D^\dagger(\bar{D})$ creates a $D(\bar{D})$ meson, τ^i is the Pauli matrix, and the ellipsis denotes other contact terms with more derivatives which are higher order terms. The first term in $\mathcal{L}_{(D\bar{D})^2}$ accounts for the interaction in the isospin singlet channel, and the second term for the isospin triplet channel. The contributions from other resonances which are not included in the Lagrangian can be saturated into the contact term $\mathcal{L}_{(D\bar{D})^2}$. Therefore, we take the coefficients such as f_1, f_3 to be complex, where the imaginary parts of them come from the width of the saturated resonances and the $D\bar{D}$ annihilation effect. With isospin symmetry, we only have to consider the terms for isospin singlet channel in $\mathcal{L}_{(D\bar{D})^2}$ to study the $D\bar{D}$ final state interactions, since the $D\bar{D}$ pair comes from the decay of $\psi(3770)$ and $\psi(2S)$ in our approach.

Now we come to the discussion on the power counting of this effective field theory. The tree level diagrams for the $D\bar{D}$ elastic scattering are shown in Fig. 1. Near the $D\bar{D}$ threshold, the denominator of the $\psi(3770)$ propagator can be expressed as

$$\begin{aligned} P(\psi) &= \frac{1}{s - M_\psi^2 + iM_\psi \Gamma_\psi^{\text{non-}D\bar{D}}} \\ &\approx \frac{1}{(2M_D + p^2/M_D)^2 - M_\psi^2 + iM_\psi \Gamma_\psi^{\text{non-}D\bar{D}}} \\ &= \frac{1}{4p^2 + 4M_D^2 - M_\psi^2 + iM_\psi \Gamma_\psi^{\text{non-}D\bar{D}} + \mathcal{O}(p^4)}, \end{aligned} \quad (4)$$

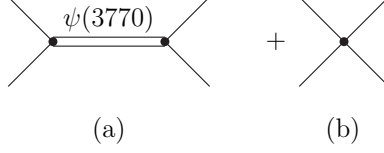


FIG. 1: Tree diagrams for $D\bar{D} \rightarrow D\bar{D}$. (a) $D\bar{D} \rightarrow \psi(3770) \rightarrow D\bar{D}$, (b) contact interaction.

where p is the magnitude of the three-vector momentum of the $D(\bar{D})$ meson in the overall center of mass frame, $\Gamma_{\psi}^{\text{non-}D\bar{D}}$ denotes the non- $D\bar{D}$ decay width of $\psi(3770)$, and M_{ψ} is the mass of $\psi(3770)$. The $D\bar{D}$ decay width of $\psi(3770)$ will be included through the summation of the D meson loops in the following. Because $\psi(3770)$ is close to the threshold of $D\bar{D}$, we expect that $P(\psi)$ is at $\mathcal{O}(p^{-2})$. Taking the momentum power of the $\psi D\bar{D}$ vertex into account, we find that Fig. 1(a) is at $\mathcal{O}(p^0)$. From the naive power counting, the leading contact terms have two derivatives, hence they are at $\mathcal{O}(p^2)$. However, in this naive power counting, we have assumed that the coefficients of the contact terms, i.e. f_1, f_3, \dots , are at order of $\mathcal{O}(p^0)$. In some cases, especially when there are bound states or resonances near threshold, the coefficients of the contact terms can be enhanced. For example, in NN interaction the S-wave contact terms C_S scale as $\mathcal{O}(p^{-1})$ [15]. Another example is the NN interaction in 3P_0 , where the leading contact term $C_{^3P_0}$ can be promoted to $\mathcal{O}(p^{-2})$ [16]. It is interesting to study whether the same enhancement mechanism takes place in the $D\bar{D}$ interactions since the resonance $\psi(3770)$ locates near the $D\bar{D}$ threshold. If f_1 is promoted to $\mathcal{O}(p^{-2})$ as $C_{^3P_0}$ in NN interaction, then the corresponding tree diagram shown in Fig. 1(b) is at $\mathcal{O}(p^0)$, which is the same as Fig. 1(a). However, since we do not know the power of f_1 at the beginning, we then assume that the leading contributions to $D\bar{D}$ elastic scattering come from both Fig. 1(a) and (b). We will use the experiment data to determine f_1 and see whether it is promoted to $\mathcal{O}(p^{-2})$. In this way, the $D\bar{D}$ scattering amplitude in the specific channel ($J^{PC} = 1^{--}, I = 0$) can be obtained by summing the ladder diagrams as shown in Fig. 2, which is equivalent to solve the Lippmann-Schwinger equation $T = V + \int VGT$ with the $D\bar{D}$ potential truncated at leading order.

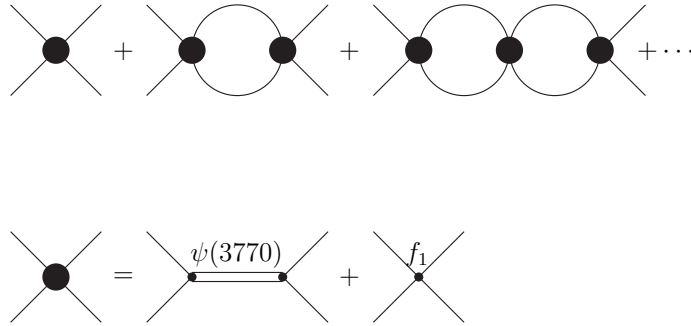


FIG. 2: The ladder diagrams for the $D\bar{D}$ interactions where the potential is truncated at the leading order.

Figure 3 illustrates the final state interactions between the produced $D\bar{D}$, and loop integrals that we will encounter in Fig. 3 can generally be reduced to

$$\begin{aligned} & \int \frac{d\ell^0 d^3\ell}{(2\pi)^4} \frac{\bar{\ell}^2}{[\ell^0 - \bar{\ell}^2/2M_D + i\epsilon] \cdot [E - \ell^0 - \bar{\ell}^2/2M_D + i\epsilon]} \\ &= -\frac{M_D}{4\pi} p^3, \end{aligned} \quad (5)$$

where $E = p^2/M_D$ is the total kinematic energy of the $D\bar{D}$ system. We can then write down the amplitude for $e^+e^- \rightarrow D\bar{D}$ as

$$i\mathcal{M} = i\mathcal{M}_a + i\mathcal{M}_b. \quad (6)$$

To be more specific, the amplitude for process Fig. 3(a) reads

$$i\mathcal{M}_a = -ie^2 \bar{v}(k_2) \gamma u(k_1) \cdot (\mathbf{p}_1 - \mathbf{p}_2) \frac{1}{s} \frac{M_\psi^2}{f_\psi} \frac{1}{s - M_\psi^2 + iM_\psi G_\psi} g_{\psi D\bar{D}}, \quad (7)$$

$$\text{with } G_\psi = \Gamma_\psi^{\text{non-}D\bar{D}} + \frac{1}{12\pi M_\psi} \left(g_{\psi D\bar{D}}^2 - f_1(s - M_\psi^2 + iM_\psi \Gamma_\psi^{\text{non-}D\bar{D}}) \right) \left(\frac{|\vec{p}_{D^0}|^3}{M_{D^0}} + \frac{|\vec{p}_{D^+}|^3}{M_{D^+}} \right), \quad (8)$$

where γ is the Dirac gamma matrix, k_1 and k_2 are the incoming momentum of electron and positron, respectively, while p_1 and p_2 are the outgoing momentum of D and \bar{D} , respectively. G_ψ cannot be simply interpreted as the width of $\psi(3770)$, since it is a complex number. If $f_1 = 0$, then $G_\psi = \Gamma_\psi^{\text{non-}D\bar{D}} + (|\vec{p}_{D^0}|^3/M_{D^0} + |\vec{p}_{D^+}|^3/M_{D^+}) g_{\psi D\bar{D}}^2/(12\pi M_\psi)$.

The amplitude for Fig. 3(b) can be written as

$$i\mathcal{M}_b = -ie^2 \bar{v}(k_2) \gamma u(k_1) \cdot (\mathbf{p}_1 - \mathbf{p}_2) \frac{1}{s} \frac{M_{\psi(2S)}^2}{f_{\psi(2S)}} \frac{1}{s - M_{\psi(2S)}^2 + iM_{\psi(2S)} \Gamma_{\psi(2S)}} \tilde{g}_{\psi(2S)}, \quad (9)$$

$$\text{with } \tilde{g}_{\psi(2S)} = g_{\psi(2S) D\bar{D}} \left[1 + i \frac{1}{12\pi} \left(-f_1 + \frac{g_{\psi D\bar{D}}^2}{s - M_\psi^2 + iM_\psi \Gamma_\psi^{\text{non-}D\bar{D}}} \right) \left(\frac{|\vec{p}_{D^0}|^3}{M_{D^0}} + \frac{|\vec{p}_{D^+}|^3}{M_{D^+}} \right) \right]^{-1}, \quad (10)$$

where $M_{\psi(2S)}$ is the mass of $\psi(2S)$ which can be read from PDG [17], and $f_{\psi(2S)}$ can be extracted by Eq. (2) using $\Gamma_{\psi(2S) \rightarrow e^+e^-} = 2.35 \text{ keV}$ [17].

To proceed, we denote the cross section for $e^+e^- \rightarrow D\bar{D}$ as $\sigma^B(s)$ which does not include the initial state radiation (ISR) effect. In reality, for a given energy \sqrt{s} , the actual c.m. energy for the e^+e^- annihilation is $\sqrt{s'} = \sqrt{s(1-x)}$ due to the ISR effect, where $x E_{\text{beam}}$ is the total energy of the emitted photons. To order α^2 radiative correction in the e^+e^- annihilation, the observed cross section σ^{obs} at BESII can be related to our result σ^B through [18]

$$\sigma^{\text{obs}}(s) = (1 + \delta_{VP}) \int_0^{1-4M_D^2/s} dx f(x, s) \sigma^B(s(1-x)), \quad (11)$$

where $(1 + \delta_{VP}) = 1.047$, and the function $f(x, s)$ is given by

$$\begin{aligned} f(x, s) &= \beta x^{\beta-1} \delta^{V+S} + \delta^H, \\ \beta &= \frac{2\alpha}{\pi} (\ln \frac{s}{m_e^2} - 1), \\ \delta^{V+S} &= 1 + \frac{3}{4}\beta + \frac{\alpha}{\pi} \left(\frac{\pi^2}{3} - \frac{1}{2} \right) + \frac{\beta^2}{24} \left(\frac{1}{3} \ln \frac{s}{m_e^2} + 2\pi^2 - \frac{37}{4} \right), \\ \delta^H &= -\beta \left(1 - \frac{x}{2} \right) + \frac{1}{8} \beta^2 \left[4(2-x) \ln \frac{1}{x} - \frac{1+3(1-x)^2}{x} \ln(1-x) - 6 + x \right]. \end{aligned} \quad (12)$$

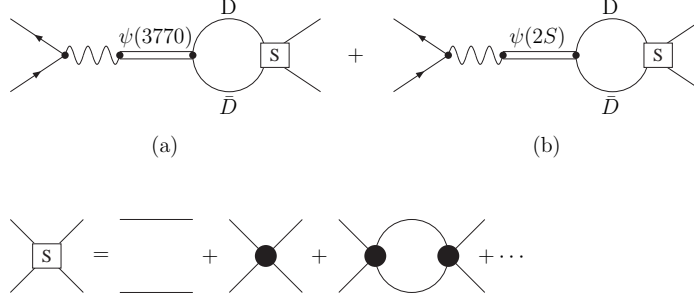


FIG. 3: The Feynman diagrams for $e^+e^- \rightarrow D\bar{D}$ in our approach.

Before fitting the BESII data with Eq. (11), we first discuss our treatment of $\Gamma_{\psi}^{\text{non-}D\bar{D}}$. It seems impossible to determine $\Gamma_{\psi}^{\text{non-}D\bar{D}}$ definitely in our fitting since $\Gamma_{\psi}^{\text{non-}D\bar{D}}$ is always accompanied by f_1 in our formula, and any change of $\Gamma_{\psi}^{\text{non-}D\bar{D}}$ can be compensated by the tuning of f_1 . The experimental results on non- $D\bar{D}$ branching ratio of $\psi(3770)$ decay are still controversial [19–21]. In contrast, the next-to-leading-order (NLO) pQCD calculation expects the non- $D\bar{D}$ decay branching ratio to be at most about 5% [22], while an effective Lagrangian approach estimates that the D meson loop rescatterings into non- $D\bar{D}$ light vector and pseudoscalar mesons leads to about 1% non- $D\bar{D}$ branching ratios [23]. A similar calculation by Ref. [24] also confirms such a nonperturbative phenomenon. Another experimental observation is that so far most of the well-measured non- $D\bar{D}$ decay modes of $\psi(3770)$ are found rather small, namely, their branching ratios are either at the order of $10^{-3} - 10^{-4}$, or only an upper limit is set [17].

Taking all these facts into account and for the purpose of studying the dominant $D\bar{D}$ channel, we set $\Gamma_{\psi}^{\text{non-}D\bar{D}}$ to be zero in our fitting as a leading approximation. We have checked that the fitting results are almost unchanged even though we set the non- $D\bar{D}$ breaching ratio of $\psi(3770)$ decay to be at the order of several percents. Our fitting results for the $D\bar{D}$ cross section lineshape are presented in Fig. 4. The fitted parameters are

$$\begin{aligned} M_{\psi} &= 3.7674(44) \text{ GeV}, & \frac{e}{f_{\psi}} &= 4.25(13) \times 10^{-3}, \\ g_{\psi D\bar{D}} &= 15.43 \pm 2.7, & g_{\psi(2S) D\bar{D}} &= -6.9 \pm 3.6, \\ f_1 &= 2059.2 \pm 534 \text{ (GeV}^{-2}\text{)}. \end{aligned} \quad (13)$$

The fitting quality is $\chi^2/\text{d.o.f.} = 23.44/22 \approx 1.07$. The imaginary part of f_1 from the fitting is almost zero, which is consistent with the fact that the $D\bar{D}$ annihilation will be highly suppressed in the heavy quark limit. It seems that the magnitude of f_1 is large, and it is worth noting that this value is at the order of $(M_D/\delta)^2$, where $\delta = M_{\psi(3770)} - 2M_D$. Therefore, the large f_1 is consistent with our previous assumption that f_1 will be promoted by the low energy scale δ , up to $\mathcal{O}(p^{-2})$. With the fitting parameters,

we can obtain

$$\begin{aligned}\Gamma_\psi &= \frac{g_{\psi D\bar{D}}^2}{48\pi M_\psi^2} \left[(M_\psi^2 - 4M_{D^0}^2)^{3/2} + (M_\psi^2 - 4M_{D^+}^2)^{3/2} \right] = 27.58 \pm 11.07 \text{ MeV}, \\ \Gamma_{ee} &= \frac{1}{3} \alpha M_\psi \left(\frac{e}{f_\psi} \right)^2 = 165.6 \pm 10.1 \text{ eV}.\end{aligned}\tag{14}$$

The values of M_ψ and Γ_{ee} we obtain are somewhat smaller than those given by PDG2010 [17], while Γ_ψ is consistent with the PDG value [17]. We note that the $\psi(3770)$ resonance parameters, such as M_ψ , Γ_ψ and Γ_{ee} , are still under improvement as shown by the PDG2012 revisions [1].

In Fig. 5, we also present the Born cross section $\sigma^B(s)$ for $e^+e^- \rightarrow D\bar{D}$ which is denoted by the solid red curve. The dashed and dotted lines are for the neutral and charged D meson pair Born cross sections, respectively. Combining the results shown in Figs. 4 and 5, we find that the anomalous cross section lineshape could be originated from the interferences from the $\psi(2S)$ pole and $D\bar{D}$ final state interactions. Since the $\psi(2S)$ pole is relatively isolated due to its relatively narrow width in comparison with the mass gap between $\psi(2S)$ and $\psi(3770)$, the relative phase between the $\psi(2S)$ and $\psi(3770)$ amplitudes is likely to be produced by the $D\bar{D}$ final state interactions. Although our calculation cannot determine the absolute value for the possible non- $D\bar{D}$ decay branching ratio of $\psi(3770)$, it is constructive to recognize the important role played by the final state $D\bar{D}$ interactions that cause the deviation of the $e^+e^- \rightarrow D\bar{D}$ cross section at the $\psi(3770)$ mass region from a Breit-Wigner shape. This analysis is useful for our further understanding of the $\psi(3770)$ non- $D\bar{D}$ decays as a manifestation of possible nonperturbative QCD mechanisms.

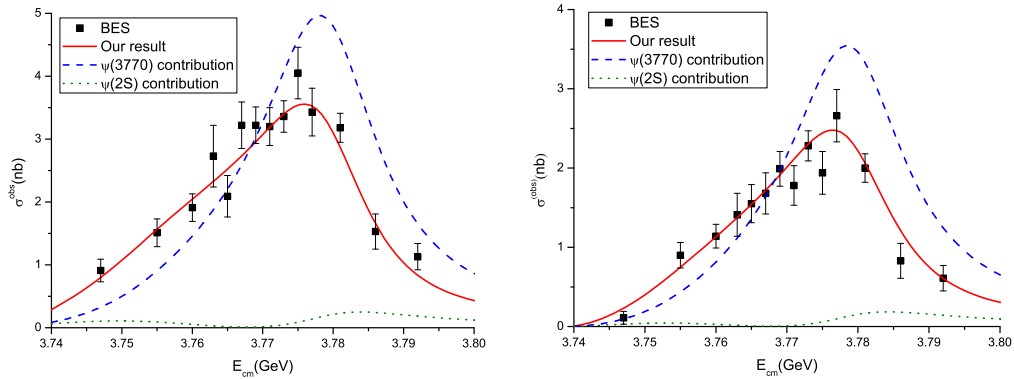


FIG. 4: The observed cross sections for $e^+e^- \rightarrow D^0\bar{D}^0$ (left plot) and $e^+e^- \rightarrow D^+D^-$ (right plot). The solid line is the fitting result in our approach shown in Fig.3, the dashed line shows the contribution from Fig.3.(a), and the dotted line shows the contribution from Fig.3.(b). The data are from BES[2].

In summary, we have proposed an effective field theory for low energy $D\bar{D}$ interactions, in which we have included the resonance $\psi(3770)$ and an additional small scale δ . It is found that the small scale δ will promote the coefficient of the contact term f_1 to be $\mathcal{O}(p^{-2})$. So the leading $D\bar{D}$ interaction potential in this specific channel would come from the S -channel $\psi(3770)$ exchange and the contact term f_1 . With the leading $D\bar{D}$ potential, we then sum the ladder diagrams to describe the $D\bar{D}$ final state interaction as shown in Fig. 3. We find that we can describe the anomalous cross section lineshape of $e^+e^- \rightarrow D\bar{D}$ observed by the BESII Collaboration [2] using the effective field theory. This approach should be useful for our further understanding of the $\psi(3770)$ non- $D\bar{D}$ decays which could be sharing the same dynamic origin as the $D\bar{D}$ cross section lineshape anomaly as emphasized in Refs. [23, 25].

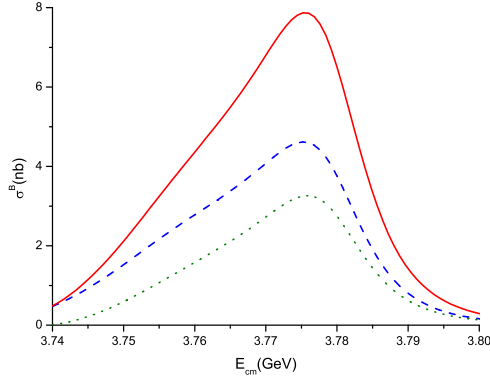


FIG. 5: The Born cross section $\sigma^B(s)$ for $e^+e^- \rightarrow D\bar{D}$. The solid line is for $e^+e^- \rightarrow D\bar{D}$, the dashed line is for $e^+e^- \rightarrow D^0\bar{D}^0$, and the dotted line is for $e^+e^- \rightarrow D^+D^-$.

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